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Mechanism of Dissipation in Heavy-Ion Reactions

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ABSTRACT

We discuss a new surface-plus-window mechanism for the conversion of nuclear collective energy into internal degrees of freedom at intermediate excitation energies. This novel dissipation mechanism, which results from the long mean free path of nucleons inside a nucleus, involves interactions of either one or two nucleons with the moving nuclear surface and also, for dumbbell-like shapes encountered in heavy-ion reactions and fission, the transfer of nucleons through the window separating the two portions of the system. To illustrate the effect of surface-plus-window dissipation on heavy-ion fusion reactions we present dynamical calculations for values of the dissipation strength corresponding to 27% and 100% of the Swiatecki wall-formula value, as well as for no dissipation. In addition to dynamical thresholds for compound-nucleus formation in heavy-ion reactions, our new picture describes such other phenomena as experimental mean fission-fragment kinetic energies for the fission of nuclei throughout the periodic system, enhancement in neutron emission prior to fission, short scission-to-scission times in sequential ternary fission, widths of mass and charge distributions in deep-inelastic heavy-ion reactions, and widths of isoscalar giant quadrupole and giant octupole resonances.

1. INTRODUCTION

Nuclear physicists have been trying for years to determine the mechanism and magnitude of nuclear dissipation in large-amplitude collective motion such as occurs in heavy-ion reactions and fission. Up until about 1974 it was generally believed that the mechanism of nuclear dissipation is two-body collisions, like that responsible for ordinary viscosity in fluids, and that the magnitude is sufficiently small that nuclei are mobile, like mercury.[1] But then, through the work of Gross, Swiatecki, and others, it was realized that the long mean free path of nucleons inside a nucleus, arising from the Pauli exclusion principle for fermions, alters both the mechanism and magnitude.[2,3] However, the subsequent approximations made in incorporating this single physical principle have led to radically different pictures.

By assuming that the velocity distribution of nucleons striking a moving container wall is completely random, Swiatecki and his colleagues derived a simple wall formula for describing such one-body dissipation, in terms of which nuclei are predicted to be superviscid.[3-8] In contrast, by constraining the many-body wave function to at all times be a Slater determinant of single-particle wave functions, Bonche, Davies, Koonin, Negele, and others treated the dynamics by use of the time-dependent Hartree-Fock approximation, in terms of which nuclei dissipate much less energy.[9-11] Attempts to experimentally discriminate between such possibilities have thus far proved elusive because of the difficulty of distinguishing dissipative effects from analogous effects caused by collective degrees of freedom.

2. SURFACE-PLUS-WINDOW DISSIPATION

We present here a new macroscopic approach to this problem, valid for intermediate excitation energies above which pairing has disappeared and below which the nucleon mean free path exceeds the nuclear diameter. For such excitation energies, the dissipation proceeds primarily in the surface region from two distinct mechanisms. The first mechanism is one-body dissipation, but with a magnitude that is substantially reduced relative to that of the wall formula. In calculations based on the random-phase approximation for spherical nuclei, Griffin, Dworzecka, and Yannouleas have shown that the effect of replacing three idealizations of the wall formula by more realistic features appropriate to real nuclei is to reduce the one-body dissipation coefficient to roughly 10% of the wall-formula value.[12,13] Alternatively, the reduction could arise because the nucleons retain some memory of their previous collisions with the wall, which invalidates the assumption of a random velocity distribution that was used to derive the wall formula. The second mechanism is two-body collisions in the surface region. The Pauli exclusion principle, which suppresses two-body collisions in the nuclear interior, disappears as one passes through the nuclear surface to the exterior.[14] In addition, the free nucleon-nucleon cross section itself increases as one passes through the surface to the exterior because of its increase with decreasing kinetic energy. Since the density decreases to zero outside the nucleus, the probability for two-body collisions peaks in the nuclear surface.

We assume that the surface dissipation is local and calculate it from the leading term in an expansion of the time rate of change of the collective Hamiltonian H in powers of the surface diffuseness divided by the nuclear radius.[5] We write this leading term as

$$\left(\frac{dH}{dt}\right)_{\text{surface}} = -k_s \rho \bar{v} \int (\dot{n} - D)^2 dS \quad , \tag{1}$$

where $\dot{n}$ is the velocity of a surface element dS, D is the normal drift velocity of nucleons about to strike the surface element dS, $\bar{v}$ is the average speed of the nucleons inside the nucleus, ρ is the nuclear mass density, and k_s is a dimensionless parameter that specifies the total strength of the interaction of either one or two nucleons with the moving nuclear surface. A value of $k_s = 1$ would correspond to the wall formula, but several types of experimental data indicate that for real nuclei its value is much less than unity. The value of k_s could depend upon both the excitation energy and type of collective motion involved.

For dumbbell-like shapes, the transfer of nucleons through the window separating the two portions of the system leads to an additional dissipation that is analogous to the classical window formula of Swiatecki.[3-8] Our result is

$$\left(\frac{dH}{dt}\right)_{\text{window}} = -\frac{1}{2}\rho \bar{v} a \dot{r}^2 F(q, \dot{q}) \quad , \tag{2}$$

where a is the area of the window, $\dot{r}$ is the relative velocity of the centers of mass of the two portions of the system, and $F(q, \dot{q})$ describes the effect of a nonuniform velocity as a function of position in the deforming fragments. There is no need to renormalize this part of the dissipation because nucleons that have passed through a small window have a low probability of returning through it while still retaining memory of their previous passage.

The combination of these two mechanisms leads to surface-plus-window dissipation. In calculating the total dissipation rate dH/dt, we describe the transition from the pure surface dissipation that applies to mononuclear shapes, where the drift D is zero in Eq. (1), to the surface-plus-window dissipation that applies to dinuclear shapes, where the drift D is nonzero in Eq. (1), by use of a smooth interpolation analogous to that used in Ref. 15.

3. MACROSCOPIC-MICROSCOPIC METHOD

To describe the axially symmetric nuclear shapes that are considered, we use the cylindrical Legendre-polynomial expansion of Trentalange *et al.*,[16] with five independent symmetric and five independent asymmetric shape coordinates. In addition, we include an angular coordinate to describe the rotation of the nuclear symmetry axis in the reaction plane, which leads to a total of 11 collective coordinates q.

We consider excitation energies that are sufficiently high that single-particle effects may be neglected and calculate the potential energy of deformation $V(q)$ as the sum of repulsive Coulomb and centrifugal energies and an attractive Yukawa-plus-exponential potential,[17] with constants determined in a recent nuclear mass formula.[18] This generalized surface energy takes into account the reduction in energy arising from the nonzero range of the nuclear force in such a way that saturation is ensured when two semi-infinite slabs are brought into contact.

The collective kinetic energy is a quadratic function of the collective velocities $\dot{q}$, with coefficients that are related to the elements of the shape-dependent inertia tensor $M(q)$. At the high excitation energies and large deformations considered here, where pairing correlations have disappeared and near crossings of single-particle levels have become less frequent, the rotational moment of inertia is close to the rigid-body value and the vibrational inertia is close to the incompressible, irrotational value.[19] We therefore calculate the inertia tensor $M(q)$, which is a function of the shape of the system, for a superposition of rigid-body rotation and incompressible, nearly irrotational flow. For this purpose we use the Werner-Wheeler method, which determines the flow in terms of circular layers of fluid.[20]

The coupling between the collective and internal degrees of freedom gives rise to dissipative forces, whose mean values are calculated for our new surface-plus-window model that was described in Sec. 2. The residual fluctuating forces are treated under the Markovian assumption that they do not depend upon the system's previous history. At high excitation energies, where classical statistical mechanics is valid, this leads to a generalized Fokker-Planck equation for the dependence upon time t of the distribution function $f(q, p, t)$ in phase space of collective coordinates and momenta.[15] In most of our studies here we use equations for the time rate of change of the first moments of the distribution function, with the neglect of higher moments. These are the generalized Hamilton equations,[15] which we solve numerically for each of the 11 generalized coordinates and momenta.

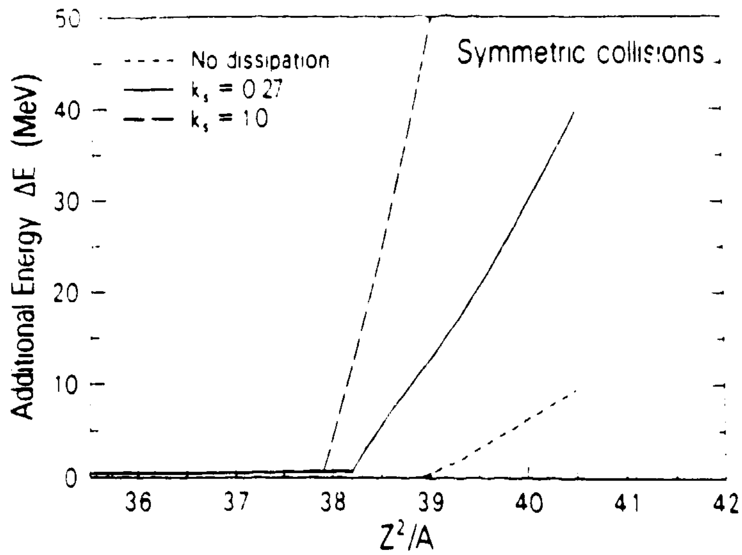

Fig. 1. Effect of dissipation on the additional bombarding energy ΔE required to form a compound nucleus.

4. HEAVY-ION REACTIONS

A necessary condition for forming a compound nucleus in a heavy-ion reaction is that the dynamical trajectory for the fusing system pass inside the fission saddle point in a multidimensional space. The dynamical trajectories and hence the cross sections for forming a compound nucleus depend strongly upon the location of the fission saddle point relative to the contact point. For light nuclear systems the fission saddle point lies outside the contact point, so that all of the angular-momentum states that cross the one-dimensional interaction barrier automatically pass inside the saddle point. However, for heavy nuclear systems and/or large impact parameters, the fission saddle point lies inside the contact point, and the center-of-mass bombarding energy must exceed the maximum in the one-dimensional zero-angular-momentum interaction barrier by an amount ΔE in order to form a compound nucleus.

This additional bombarding energy ΔE has been calculated over the past 14 years by use of various approximations and for several dissipation mechanisms.[6-8,15,21-26] Here we concentrate on results that we have obtained recently by solving the generalized Hamilton equations numerically for our new surface-plus-window mechanism, with illustrative values of $k_s = 0.27$ and 1.0. The former value, which was initially determined from an adjustment to experimental isoscalar giant quadrupole and giant octupole widths,[25] lies within the limits $0.2 \leq k_s \leq 0.5$ determined from a comparison of calculated and experimental mean fission-fragment kinetic energies,[26] whereas the latter value corresponds to the Swiatecki wall formula.[3-8] To provide reference points, we in all cases perform similar calculations for no dissipation. The results presented here are all calculated for systems in which the target and projectile each lie along Green's approximation to the valley of β stability.[27]

The dependence of the additional energy ΔE upon the size of the compound nucleus that is formed in symmetric collisions is shown in Fig. 1. For no dissipation, ΔE is zero below the threshold value $(Z^2/A)_{thr} = 38.9$ and then increases slowly with increasing Z^2/A. With dissipation, ΔE is slightly positive for all values of Z^2/A because of the

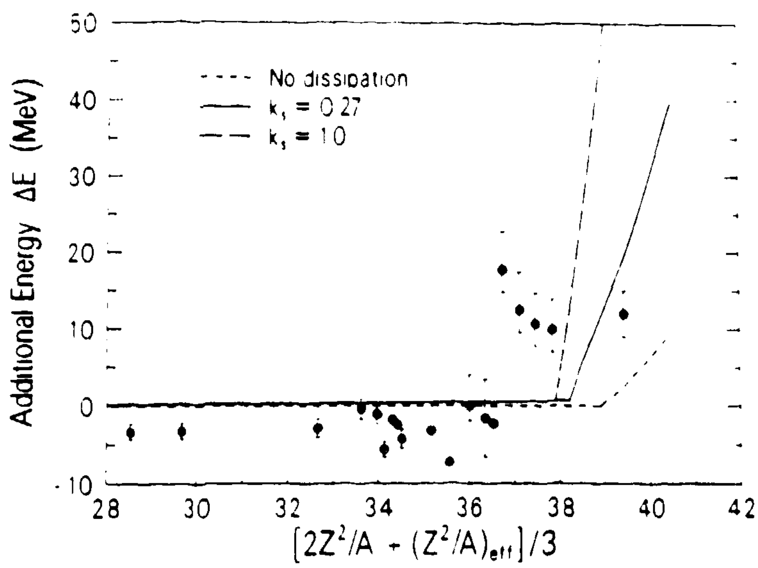

Fig. 2. Comparison of calculated and experimental values of the additional bombarding energy ΔE required to form a compound nucleus.

energy dissipated during the approach of the target and projectile prior to their reaching the maximum in the one-dimensional zero-angular-momentum barrier. For surface-plus-window dissipation with $k_s = 0.27$, the threshold value is lowered to $(Z^2/A)_{thr} = 38.2$, above which ΔE increases rapidly with increasing Z^2/A. Finally, for $k_s = 1.0$, the threshold value is lowered even further to $(Z^2/A)_{thr} = 37.9$, above which ΔE increases very rapidly with increasing Z^2/A.

The large differences between the three curves in Fig. 1 offer the tantalyzing possibility of determining the mechanism and magnitude of nuclear dissipation from comparisons with experimental data. However, because most of the systems that are studied experimentally are asymmetric, it is necessary to either perform calculations for many individual asymmetric systems or to scale the results for asymmetric systems into those calculated for symmetric systems. For this purpose we use the weighted average $(Z^2/A)_{wt}$ defined[8,26] by the abscissa in Fig. 2, where

$$(Z^2/A)_{eff} = \frac{4Z_1Z_2}{A_1{}^{1/3}A_2{}^{1/3}(A_1{}^{1/3} + A_2{}^{1/3})} \quad . \tag{3}$$

With this choice of scaling abscissa, we compare in Fig. 2 our results calculated for symmetric systems with experimental values of ΔE derived from measurements of evaporation residues for various asymmetric systems.[8] Unlike in the comparisons made in Ref. 8, where the experimental values of ΔE were measured relative to barrier heights calculated by use of the Bass potential,[28] in Fig. 2 the experimental and calculated values of ΔE are both measured relative to barrier heights calculated by use of the Yukawa-plus-exponential potential,[17] with constants determined from a recent nuclear mass formula.[18] For $(Z^2/A)_{wt} \leq 36$, the experimental values of ΔE lie systematically below all three calculated curves by an average of about 3 MeV. This discrepancy could represent an effective lowering of the experimental barrier by zero-point vibrations of the approaching target and projectile, or alternatively could represent a slight deficiency in the values of the constants that are used for the Yukawa-plus-exponential potential.

At $(Z^2/A)_{wt} \approx 37$ the experimental values of ΔE show a rapid rise, as anticipated. However, on a finer scale the trend of the experimental results with $(Z^2/A)_{wt}$ does not follow that predicted by any of the curves. The decrease of ΔE for the four solid circles above the curves as $(Z^2/A)_{wt}$ increases between 36 and 38 is probably associated partly with a change in ground-state shape or stiffness of the $^{90,92,94,96}Zr$ targets involved. In contrast, the relatively small value of ΔE for the solid circle at $(Z^2/A)_{wt} = 39.4$ probably arises partly from a valley in the potential-energy surface that leads inward from the tangent-sphere configuration for nearly magic target and projectile. While further experimental data that differentiate between shell effects and smooth trends are clearly needed before a definitive conclusion can be drawn from such comparisons, it is apparent that our new surface-plus-window dissipation with $k_s = 0.27$ adequately describes the overall features of experimental dynamical thresholds for compound-nucleus formation.

5. OTHER PHENOMENA

The dynamical evolution of a fissioning nucleus also depends strongly on the value of the surface dissipation coefficient k_s appearing in Eq. (1). As k_s increases, the motion is slowed down and collective energy is converted into internal degrees of freedom, which reduces the mean translational kinetic energy of the fission fragments at infinity. However, this reduction is partially offset by more compact scission shapes for surface-plus-window dissipation compared to no dissipation, which increases the post-scission kinetic energy acquired through the Coulomb repulsion of the fragments. Experimental mean kinetic energies for the fission of nuclei throughout the periodic system at high excitation energy, where single-particle effects have decreased in importance, are satisfactorily reproduced for values of k_s within the range $0.2 \leq k_s \leq 0.5$.

Surface-plus-window dissipation with a relatively small coefficient is also able to describe the enhancement in neutron emission prior to fission that has been observed recently in several heavy-ion-induced reactions. For $^{16}O + ^{142}Nd \rightarrow ^{158}Er$ at a laboratory bombarding energy of 207 MeV, there are experimentally 2.7 ± 0.4 neutrons emitted prior to fission, compared to 1.6 neutrons calculated with a standard statistical model.[29] An interpretation of this enhancement in terms of neutron emission during the transient time required to build up the quasi-stationary probability flow over the barrier and the mean time required for the system to descend from the saddle point to scission yields an upper limit for the reduced nuclear dissipation coefficient[29] that is consistent with the upper limit determined for surface-plus-window dissipation from mean fission-fragment kinetic energies.[26]

The short scission-to-scission times observed experimentally in sequential ternary fission also require a relatively small dissipation coefficient. For $^{129}Xe + ^{122}Sn$ at a laboratory bombarding energy per nucleon of 12.5 MeV, where ternary events were observed approximately 10% of the time when the energy loss was high, Glässel *et al.*[30] deduced that the time between successive scission events is approximately 1×10^{-21} s. This result has been analyzed by Plasil,[31] on the basis of calculated mean saddle-to-scission times, to yield an upper limit to the dissipation strength that is also consistent with the upper limit determined for surface-plus-window dissipation.[26]

The widths of mass and charge distributions in deep-inelastic heavy-ion reactions[32] are also adequately described for small energy losses by surface-plus-window dissipation. We have performed no additional calculations in this area, but previous analyses have confirmed the correctness of the window formula.[32] This contribution is carried over essentially unchanged to our new picture, and the calculated widths are insensitive to the strength of the remaining dissipation.

As mentioned in Sec. 2, the value of k_s could depend upon both the excitation energy and type of collective motion involved. To see whether this is indeed the case, we have used our new dissipation picture to calculate the isoscalar giant-resonance width as a function of mass number A and multipole degree n, approximating such resonances by small-amplitude collective oscillations in which the neutrons and protons undergo in-phase, incompressible, irrotational flow with unit effective mass.[25,26] This type of flow pattern is expected to arise when the nucleons remain in orbitals characterized by their original nodal structure. For a value of $k_s = 0.27$, our new dissipation picture satisfactorily reproduces the average trends of the experimental widths of isoscalar giant quadrupole and giant octupole resonances as a function of mass number.[25,26]

6. SUMMARY AND CONCLUSION

Building on the best features of previous work, we have developed a new picture of nuclear dissipation that provides a unified macroscopic description of several diverse phenomena. These include dynamical thresholds for compound-nucleus formation in heavy-ion reactions, mean fission-fragment kinetic energies, enhancement in neutron emission prior to fission, short scission-to-scission times in sequential ternary fission, widths of mass and charge distributions in deep-inelastic heavy-ion reactions, and widths of isoscalar giant quadrupole and giant octupole resonances. We hope that our new picture will be subjected to countless experimental tests during the coming years.

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